

Fine-Tuning in DBI Inflationary Mechanism

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Abstract

We show a model-independent fine-tuning issue in the DBI inflationary mechanism. DBI inflation requires a warp factor h small enough to sufficiently slow down the inflaton. On the other hand, the Einstein equation in extra dimensions under the inflationary background deforms the warp space in the IR side. Generically these two locations coincide with each other, spoiling the DBI inflation. The origin and tuning of this “ h -problem” is closely related, through the AdS/CFT duality, to those of the well-known “ η -problem” in the slow-roll inflationary mechanism.

1 Introduction

One important motivation for the DBI (Dirac-Born-Infeld) inflation [1,2] is to circumvent the η -problem in the slow-roll inflation [3]. Namely, the required inflaton mass-squared for the slow-roll inflation, $\lesssim \mathcal{O}(0.01H^2)$, is generically corrected to $\mathcal{O}(H^2)$ due to the backreaction of the inflationary background. This leads to a potential too steep to support the slow-roll inflation. It is proposed that warped space can provide a speed-limit which holds the inflaton near the top of a potential even if the potential is steep. Indeed, given such a warped space, it is shown that the inflation can happen, in different situations, with a very steep potential (typically $m^2 \gg H^2$) [1] or a potential of generic shapes (typically $m^2 \sim \mathcal{O}(H^2)$) [2]. See Ref. [4] for a summary. The importance of the warped space for this mechanism is obvious.

In this paper we discuss a subtlety involved in the construction of such a warped space. At the level of the 4d effective field theory, the warping effect shows up in the kinetic term of the inflaton $r(\mathbf{x}, t)$ through the warp factor $h(r)$,

$$h^4 \sqrt{1 + h^{-4} g^{\mu\nu} \partial_\mu r \partial_\nu r} , \quad (1.1)$$

and it seems that we can take whatever $h(r)$ we like. A typical example is the AdS_5 space with a scale R ,

$$h(r) = \frac{r}{R} . \quad (1.2)$$

In order to provide a speed-limit small enough for the inflation to happen, h is required to be smaller than a critical value h_0 .

On the other hand, the warped space lies in the internal field space, or the extra dimensions. In order to have a consistent UV completion of the DBI inflation, such as in terms of brane inflation [5], the warp factor is supplied by a metric

$$ds^2 = h(r)^2 (-dt^2 + a(t)^2 d\mathbf{x}^2) + h(r)^{-2} dr^2 , \quad (1.3)$$

where $a(t)$ is the scale factor of the 4d inflationary background. This metric has to satisfy Einstein's equation in the extra dimensions. As we will see, due to the backreaction of the 4d inflationary background, the warp factor h has to be deformed away from (1.2) near a critical value. Generically, this value coincides with what is required for the inflation, h_0 , and the resulting metric leads to a huge inflaton probe backreaction. So tuning is necessary.

We comment that this tuning issue refers to the *mechanism* of the DBI inflation [1,2], which provides the bases for various observational consequences [6,7]. With this mechanism achieved, there may be other tuning issues involving model parameters when data comparison is made [6,7]. To distinguish from these secondary issues, we call the problem discussed here the “ h -problem”.

Maldacena first emphasized, a few years ago, the importance of the backreaction from the Hubble expansion to the DBI inflationary mechanism [8]. Generally we expect a warped space with a very small minimum warp factor to be deformed in the IR side by the Hubble expansion, so the location where the deformation starts to be significant is crucial. So far there have been several kinds of such backreactions discussed in the literature [7, 9–12]. Particles such as strings can be created in the IR side by the spatial expansion because their mass scales are redshifted. Only when the density of these particles exceeds certain critical value, its backreaction shortens the warped throat; but this still leaves a considerable portion of warped space for DBI inflation [7]. Various moduli that stabilize the IR warp factor in the flux compactification will fluctuate in the inflationary background and this also shortens the throat [9, 10]. The analyses along this line so far has not been explicit enough for our purpose due to technical difficulties in solving equations of motion with large warping, and also it is likely to be model-dependent. The throat can also be deformed due to the mediation of the supersymmetry breaking [11], such effects are model-dependent of the details of the supersymmetry breaking. As mentioned, in this paper we explain the most general and significant source of the backreaction to the DBI inflationary mechanism, based on the Einstein equations for the metric. The importance of identifying such a source has an analogue in the slow-roll inflation, where the most general source of the η -problem is identified to be the canonical inflaton dependence in the Kahler potential [3], and we can therefore look for other model-dependent backreactions to tune away this problem. Similarly for DBI inflation, other model-dependent backreaction aspects mentioned above, such as the particle creation, moduli shifting and the supersymmetry breaking, may provide the counter-sources to tune away the h -problem, as we will discuss.

We will also see that the origin and tuning of the h -problem in DBI inflation is closely related to those of the η -problem in slow-roll inflation through the AdS/CFT duality. Hopefully both the explicit (through the Einstein equation) and implicit (through the duality) descriptions will help illustrating the generality of this problem and toward finding the explicit solutions.

2 The h -problem

In DBI inflation, the potential is steep so the inflaton travels near the speed-limit provided by the warped space. For (1.2), the leading behavior of the inflaton is

$$r \approx \pm \frac{R^2}{t} , \quad |t| > H^{-1} . \quad (2.4)$$

In order to have DBI inflation, we need $\Delta t > H^{-1}$, so the warp factor h has to be smaller than $h_0 = HR$. This statement is model-independent once (1.2) is given. To have N_e e-folds

of inflation, for the IR DBI model [2], the warp factor has to be as small as HR/N_e where H is a constant; for the UV DBI model [1], as small as HR/p where $p \gg 1$ is a constant and H slowly drops by a factor of $e^{N_e/p}$ during the inflation.

Let us look at the Einstein equation for the warp factor. In order to have a stabilized warped compactification, we need the source for the warping as well as some bulk fields that stabilize the extra dimensions, although their roles can be mixed. The situation for a full string-theoretic construction can be quite complicated [13]. To illustrate our main point, we first use a simpler toy model, i.e. the Randall-Sundrum (RS) model [14] with the Goldberger-Wise (GW) stabilization [15, 16]. We shall turn to the type IIB string theory shortly. The IR and UV branes separated in the fifth dimension together with the bulk potential $V(\Phi)$ provide the source for the warped space; while the GW scalar field $\Phi(r)$ in the bulk stabilizes the size of the extra dimension. The Einstein equation for the warp factor $h(r)$ in (1.3) is

$$h'^2 - H^2 h^{-2} = \frac{1}{24M_5^3} \left[\frac{1}{2} h^2 \Phi'^2 - V(\Phi) \right], \quad (2.5)$$

where the Hubble parameter H is approximately a constant (which is true at least for a few e-folds in a sustained period of inflation), M_5 is the 5d Planck mass and the prime ' denotes the derivative with respect to r .

We first consider the case where the bulk field contribution on the right hand side of (2.5) to $h(r)$ is negligible and the bulk potential is a negative constant, $V = -24M_5^3/R^2$. It is not difficult to see that the term $H^2 h^{-2}$ becomes increasingly important towards the IR side, and the naive solution (1.2) is deformed near $h \approx HR$. The solution becomes

$$h = \left(\frac{r^2}{R^2} - H^2 R^2 \right)^{1/2}. \quad (2.6)$$

This solution has been worked out before [16–18], and the form we present here is related to them by simple change of variables. The singularity at the horizon $r = HR^2$ is a coordinate singularity.

Interestingly, the solution (2.6) is related to the static AdS_5 ,

$$ds^2 = \frac{\tilde{r}^2}{R^2} (-d\tilde{t}^2 + d\mathbf{x}^2) + \frac{R^2}{\tilde{r}^2} d\tilde{r}^2, \quad (2.7)$$

by a coordinate transformation [16, 19],

$$\tilde{t} = -\frac{H^{-1}e^{-Ht}}{\sqrt{1 - \frac{H^2 R^4}{r^2}}}, \quad (2.8)$$

$$\tilde{r} = R^2 H e^{Ht} \sqrt{\frac{r^2}{H^2 R^4} - 1}. \quad (2.9)$$

This observation makes it straightforward to generalize the metric (2.6) to the case of the type IIB string theory, by adding a trivial angular part $R^2 d\Omega_5^2$ and a correspondingly transformed RR 4-form [19],

$$\mathcal{C}_4 = \left(\int \frac{4}{R} h^3 dr \right) d\text{Vol}_{\text{dS}_4} . \quad (2.10)$$

At this point, any other contributions from the bulk are ignored as before.

In this case, instead of working with the naive form of the metric (1.2), the relevant region for the DBI inflation is given by (2.6). The speed-limit still approaches zero near the new horizon $r = HR^2$, but it has a very different behavior. For example, under a generic potential $V = V_0 - \frac{1}{2}\beta H^2 r^2$ ($\beta \sim 1$), the inflaton behaves as

$$\frac{r}{HR^2} - 1 = \frac{2}{e^{-2Ht} - 1} - \frac{200}{\beta^2} e^{6Ht} + \dots , \quad t < -H^{-1} , \quad (2.11)$$

where the first term gives the speed of light, and the second term is determined by the RR-potential (2.10) and V .¹ Notice that the inflaton now approaches the horizon exponentially fast with $|t|$, in contrast to the inverse relation in Eq. (2.4). This leads to the rapid exponential growth of the Lorentz factor

$$\gamma \approx \frac{\beta}{20} e^{2N_e} , \quad (2.13)$$

where $N_e \approx H|t|$ is the number of e-folds to the end of inflation. This is in contrast to the much milder growth behavior (such as linear [2]) in the AdS geometry (1.2). Now the main problem is the rapid growth of the probe backreaction. In terms of brane inflation in warped compactification, to support an ultra-relativistic probe brane with Lorentz factor γ , the warped throat needs to have a charge $\gg \gamma$ [1, 20]. To have 60 e-folds, the growth in (2.13) is undesirably large.²

¹ To derive this solution from the equation of motion

$$\frac{d}{dt} \left(\frac{h^2 \dot{r}}{\sqrt{h^4 - \dot{r}^2}} \right) + 3H \frac{h^2 \dot{r}}{\sqrt{h^4 - \dot{r}^2}} + \frac{\partial_r(h^2)(2h^4 - \dot{r}^2)}{\sqrt{h^4 - \dot{r}^2}} - \partial_r C_4 + \partial_r V = 0 , \quad (2.12)$$

one can expand $r(t)$ around the speed of light, $r = HR^2 + 2HR^2/(e^{-2Ht} - 1) + aHR^2 e^{bHt} + \dots$, $t < -1/H$, in the region $\frac{r}{HR^2} - 1 \ll 1$ where the speed limit becomes important and where $C_4 \rightarrow \frac{16\sqrt{2}}{5} H^4 R^4 (\frac{r}{HR^2} - 1)^{5/2}$. One can see that the 2nd, 3rd and 5th terms in Eq. (2.12) are important and the coefficients a and b in the above ansatz can be determined. One can also see that the exponential dependence on t in the expansion is due to the fact that the speed of light exponentially approaches the horizon, and this is quite generic against the shape of the potential. For example adding a linear potential makes $\partial_r V$ a different constant near $r \rightarrow HR^2$, and this will only change the parameter β in (2.11).

²Because the magnitude of the non-Gaussianity is given by the estimator $f_{NL} \sim \gamma^2$, Eq. (2.13) would seem to certainly violate the observational bound. However this conclusion may be too naive since the

Since the second term on the left hand side of (2.5) is always present, we regard the above problem to be a generic feature. As we will discuss in the following sections, in order to have successful DBI inflation, other contributions such as the non-renormalizable operators or bulk fields have to be introduced to cancel this term.

The presence of the GW bulk field can also modify the term h , and therefore gives a different warping behavior from (1.2) even in the absence of the inflationary background. At least for the toy model, this does not improve the situation for a more general class of warping. Consider a different warp factor

$$h = \left(\frac{r}{R}\right)^\alpha. \quad (2.14)$$

Setting $H = 0$ one can obtain what is needed for the bulk and boundary potentials for the Φ field, by solving (2.5) and another equation of motion for Φ . We ignore such details here. The point here is that once such a warping is obtained, adding back the term with H , the deformation happens for $h'^2 < H^2 h^{-2}$, i.e.

$$h < \left(\frac{HR}{\alpha}\right)^{\frac{\alpha}{2\alpha-1}}. \quad (2.15)$$

The speed-limit of the inflaton is

$$r = \left[\pm \frac{R^{2\alpha}}{(2\alpha-1)t}\right]^{\frac{1}{2\alpha-1}}, \quad (2.16)$$

so in order to have DBI inflation,

$$h < \left(\frac{HR}{2\alpha-1}\right)^{\frac{\alpha}{2\alpha-1}}. \quad (2.17)$$

Note that the different geometry generically changes the critical value h_0 . However the position of the backreaction deformation is also changed. To have (2.15) much smaller than (2.17), we need α to be very close to $1/2$. But in such a case the speed of light approaches the horizon near exponentially with t , and this will make the inflaton Lorenz factor grow rapidly (close to exponentially) and lead to a very large probe backreaction as we considered previously.

3 Relation to the η -problem

In the spirit of Maldacena's conjecture [21], in the type IIB string theory, a warped space with $\text{AdS}_5 \times S^5$ attached to a stabilized six dimensional bulk is dual to a strongly coupled $\mathcal{N} = 4$

stringy effects are rapidly becoming important due to the warping and the relativistic effect, so that the field-theoretic results for the primordial fluctuations may not apply in the relevant window [4, 7]. Here we stick to the model consistency conditions before any comparison with observational data is made.

SYM field theory coupled to gravity. Aspects of this duality are discussed in Ref. [22, 23]. For our interest, we impose an inflationary background on both sides of the duality and study the properties of fields in this background. We study a dilaton field ϕ in the deformed AdS space in the higher dimensional string theory and then inspect what it means for the dual lower dimensional field theory.

The equation of motion for the s-wave dilaton,

$$\frac{1}{\sqrt{-G}}\partial_M(\sqrt{-G}G^{MN}\partial_N\phi) = 0 , \quad (3.18)$$

becomes

$$\frac{\partial^2\phi}{\partial t^2} + 3H\frac{\partial\phi}{\partial t} - \frac{\nabla^2}{a^2}\phi - \frac{1}{h}\frac{\partial}{\partial r}\left(h^5\frac{\partial}{\partial r}\phi\right) = 0 . \quad (3.19)$$

Separating the variables,

$$\phi(t, r) = \varphi(r)g(\mathbf{x}, t) , \quad (3.20)$$

we get two equations,

$$\frac{1}{h}\frac{d}{dr}\left(h^5\frac{d}{dr}\varphi\right) + m^2\varphi = 0 , \quad (3.21)$$

$$\ddot{g} + 3H\dot{g} - \frac{\nabla^2}{a^2}g + m^2g = 0 , \quad (3.22)$$

where m^2 is the eigenvalue of the first differential equation. Redefining

$$y^2 \equiv \frac{r^2}{H^2 R^4} - 1 , \quad (3.23)$$

Eq. (3.21) has only one dimensionless parameter m^2/H^2 ,

$$y^2(1+y^2)\frac{d^2\varphi}{dy^2} + y(5y^2+4)\frac{d\varphi}{dy} + \frac{m^2}{H^2}\varphi = 0 . \quad (3.24)$$

Starting with the normalizable solution $\varphi \sim 1/y^4$ at $y \gg 1$, approaching $y \ll 1$, φ generally behaves as

$$\varphi = c_+ y^{n_+} + c_- y^{n_-} . \quad (3.25)$$

The two linearly independent solutions are given by

$$n_{\pm} = -\frac{3}{2} \pm \sqrt{\frac{9}{4} - \frac{m^2}{H^2}} , \quad (3.26)$$

where n_{\pm} can be complex (φ is real). First we can see that $m^2/H^2 \leq 0$ is unacceptable. Multiplying Eq. (3.24) by $y^2\varphi/\sqrt{1+y^2}$ and integrating over y , we get

$$\left[y^4 \sqrt{1+y^2} \varphi \frac{d\varphi}{dy} \right]_{y_m}^{\infty} + \int_{y_m}^{\infty} dy \left[\frac{m^2}{H^2} \frac{y^2}{\sqrt{1+y^2}} \varphi^2 - y^4 \sqrt{1+y^2} \left(\frac{d\varphi}{dy} \right)^2 \right] = 0, \quad (3.27)$$

where we denote the lower limit as y_m . For the normalizable and regular solution, the first term above vanishes and this equation cannot be satisfied if $m^2 \leq 0$.

If we trust (2.6) all the way to $r = HR^2$, both solutions are irregular at the horizon.³ In fact there should be a smoothing cutoff at a small y_m , because the local blue-shifted Hubble parameter is infinite at $y = 0$ and particle creation will become important. The regular solution is possible with this cutoff by requiring $d\varphi/dy|_{y_m} = 0$. This can be satisfied for $m^2/H^2 > 9/4$ where the solution starts to oscillate. The periodicity of the trigonometric function determines the spacing between the discrete eigenvalues to be roughly $\Delta\sqrt{m^2/H^2 - 9/4} \sim \pi/|\ln y_m|$. Indeed, numerical calculation finds, for example for $y_m = 0.01$, that $m^2/H^2 = 2.683, 3.896, 5.825, \dots$. So we have shown that the eigenvalue m^2 has a positive gap of order H^2 .

Let us look at the second equation (3.22). This is the familiar equation of motion for a scalar field with mass m in a dS_4 background. Now we can come back to the duality conjecture mentioned at the beginning of this section. The dual 4d particle state can be thought of as having a normalizable wave-function in the direction of the AdS space from the dual higher dimensional point of view. The specific deformation of this wave-function at the IR end of the AdS space is translated into an effective mass gap of order H from the 4d field theory point of view. Qualitatively the appearance of this mass gap is the origin of the field theoretic η -problem in the slow-roll inflation. For any candidate scalar inflaton, such a mass gives contribution of order unity to the slow-roll parameter $\eta = M_{\text{Pl}}^2 V''/V$.

The development of the mass gap is due to the IR cutoff near the horizon $r = HR^2$. Otherwise, the IR direction of the AdS space is effectively non-compact for the wave-function, and the 4d field theory is scale-invariant. This is qualitatively similar to what happens in the confinement and finite temperature theories [24]. This is also qualitatively similar to the general argument of the mass spectrum of the KK particles in such a warped space with IR cutoff at h_0 , i.e. the spectrum is discretized in the unit h_0/R where h_0 is now HR .

For a comparison, Ref. [11] (P26) discussed a fine-tuning issue that has some interesting relation to the arguments in this section, but with a different physical origin. There one starts

³ In this case, even if we only look for normalizable solutions, we still cannot find the eigenvalue due to the singular behavior of the solutions at the horizon. To be normalizable near $y = 0$, the metric (2.6) requires $\varphi \sim y^n$ with $n > -3/2$. This leaves a possible region $0 < m^2/H^2 < 9/4$ for the branch $\varphi \sim y^{n+}$. However numerical calculation suggests no eigenvalue due to the rapid growth (as $y \rightarrow 0$) of the other non-normalizable branch.

by assuming a mass gap $m \sim F/M_{\text{Pl}}$ through gravity mediation (F is the supersymmetry breaking F-term) on the field theory side, which in principle does not have to be the same as our H but may be regarded analogously as our η -problem. Then this mass gap should be dual to a cutoff in the AdS geometry by running our argument reversely (although without an explicit metric). The starting mass gap is in the strongly coupled field theory side and is not explicit so far. In addition, Ref. [11] uses a chaotic potential with the same mass $m \sim F/M_{\text{Pl}}$ for the probe brane through gravity mediation. This is a special case, since for the chaotic potential $m^2\phi^2$ the DBI inflation requires $m \gg M_{\text{Pl}}/\sqrt{N}$ (where N is the effective throat charge) [1], hence no specific relation to m defined above. The consistency criteria used to see the problem is that the potential energy dominates over the kinetic energy. Although for the UV model, this leads to a tuning problem, it will be a very weak constraint for the IR DBI model (since the potential there is dominated by a constant) and therefore will not lead to a fine-tuning problem. In this paper, we work on the string theory side first by identifying the source term of the backreaction in the Einstein equation and examining the explicit form of the deformed metric. The criteria used is the universal speed-limit (2.4). We saw that the source of the h -problem is solely due to the inflationary background with the Hubble parameter H , and is independent of specific models and of whether the probe brane is coupled, or the mass gap is generated, through the gravity or gauge mediation. (In our point of view, such model dependent aspects in Ref. [11] may actually become important as the counter-sources used to tune away the h -problem discussed here.) So the fine-tuning issue in the h -problem is more explicit and model-independent. Here the corresponding mass gap in the field theory side is only implied via the duality.

There is also another aspect in which the h and η -problems are related. We introduce a probe brane, namely a candidate inflaton, in this geometry. We have seen in the previous section that the DBI inflation is spoiled. We now consider the probe brane moving much slower than the speed-limit and see if it can satisfy the slow-roll condition. In this case, one can perturbatively expand the DBI action (1.1). The interesting region is now $r \gg HR^2$. Although in the leading order the deformed gravity field (2.6) and the deformed RR field (2.10) cancel, at the subleading order they do not. The net contribution is a mass term for the probe brane, $m^2 = H^2$ [25]. This should be viewed as the dual string theory description [25] of the conformal coupling in the field theory description of the KKLMMT slow-roll inflation [26]. So the same deformed geometry causes both the h -problem and η -problem, this time both in the string theory side.

4 Tuning

The full deformation of the metric due to the inflation should be more subtle than that due to the AdS blackhole. In light of the AdS/CFT duality relation discussed above, we can think of this problem in the 4d field theory side. In the AdS blackhole case [24], the dual field theory has a finite temperature T . One can fix T while sending the cutoff scale to infinity. In other words, the correction to the mass gap from the non-renormalizable operators, i.e. those irrelevant operators suppressed by the UV cutoff scale such as M_{Pl} , will be suppressed by M_{Pl} . In the inflationary case we cannot do that since H itself is suppressed by the Planck mass M_{Pl} (typically with the inflationary energy scale or supersymmetry breaking scale holding fixed). Therefore on the field theory side various non-renormalizable operators may provide an effective mass-squared of the same order of magnitude H^2 , but with an opposite sign. From the discussion below Eq. (2.4), a reduction of the mass gap from the order H to $0.01H$ on the field theory side will open up a portion of warped space on the string theory side for 100 e-folds of DBI inflation. This is interestingly similar amount of tuning as for the slow-roll potential. In terms of the warped compactification, this tuning depends on the details of the UV completion, namely the way in which the warped space is attached to the bulk and the physics in the bulk such as the supersymmetry breaking. For example, such non-renormalizable operators are also expected to be the generic tuning source for the slow-roll inflationary potential. Constructing explicit examples will be a very interesting open question.

Another related way to provide the tuning is to arrange the bulk field contribution on the right hand side of the equation of motion (2.5) to cancel the second term with H on the left hand side. This amounts to for example setting up the potentials for the bulk field. This approach is more feasible in more complicated examples instead of this toy model.

To conclude, we have shown that, in the DBI inflationary mechanism, there is a fine-tuning h -problem concerning the explicit construction of the warped space taking into account of the backreaction from the inflationary background. The non-renormalizable operators in the dual field theory and the bulk fields should play important roles in such fine-tuning. Without tuning, the part of the naive warped space required for the DBI inflation will be deformed by the inflationary background through Einstein's equation in the extra dimensions. In the deformed geometry, the probe backreaction of the inflaton will spoil the DBI inflation unless the background throat charge is undesirably large. This resembles the familiar fine-tuning η -problem in the slow-roll inflationary mechanism concerning the explicit construction of the flat potential taking into account of the backreaction from the inflationary background. The origins and tunings of these two problems are interestingly related through the AdS/CFT relation. Nonetheless, these two inflationary mechanisms have

distinctive observational predictions, and should reveal different aspects of the underlying fundamental theory [4, 27–32].

Acknowledgments

I would like to thank Allan Adams, Andrew Frey, Shamit Kachru, Justin Khoury, Louis Leblond, John McGreevy, Wati Taylor, and especially Hong Liu, Juan Maldacena, Fernando Quevedo for helpful discussions, and Eva Silverstein for explaining the discussions in Ref. [11] related to our Sec.3. This work is supported by the US Department of Energy under cooperative research agreement DEFG02-05ER41360.

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